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The Decay $b \rightarrow s\gamma$, the Higgs Boson Mass, and Yukawa Unification without R-Parity¹

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Abstract

We review some properties of Bilinear R-Parity Violating models: simple extensions of the Minimal Supersymmetric Standard Model motivated by spontaneous breaking of R-Parity. We concentrate on the relaxation of the bounds on the charged Higgs mass imposed by the measurement of $B(b \rightarrow s\gamma)$, the effect on the mass of the lightest neutral Higgs boson, the radiative breaking of the electroweak symmetry, the unification of bottom and tau Yukawa coupling, and the relation of these phenomena to the radiatively generated tau-neutrino mass.

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1 Introduction

Explicit Bilinear R–Parity Violation (BRpV) is a simple extension of the Minimal Supersymmetric Standard Model (MSSM) which has the attractive feature of generating neutrino masses radiatively, thus naturally small. The origin of the neutrino mass is then linked to supersymmetry [1] through the mixing of neutral higgsinos and gauginos with the neutrino. The study of models which include BRpV terms, and not trilinear (TRpV), is motivated by spontaneous R–Parity breaking [2], where only BRpV terms are generated in the superpotential. The simplest model includes violation of R–Parity and lepton number only in the tau sector, and the superpotential is

$$W = h_t \hat{Q}_3 \hat{U}_3 \hat{H}_2 + h_b \hat{Q}_3 \hat{D}_3 \hat{H}_1 + h_\tau \hat{L}_3 \hat{R}_3 \hat{H}_1 - \mu \hat{H}_1 \hat{H}_2 + \epsilon_3 \hat{L}_3 \hat{H}_2, \quad (1)$$

where ϵ_3 has units of mass and is the only extra term compared with the MSSM. The presence of this term induces a non–zero vacuum expectation value v_3 of the tau–sneutrino [3, 4, 5]. The study of different aspects of the phenomenology of this model may be simpler in different basis [6]. Besides the original one in eq. (1), a useful basis is defined by the rotation $\mu' \hat{H}'_1 = \mu \hat{H}_1 - \epsilon_3 \hat{L}_3$ and $\mu' \hat{L}'_3 = \epsilon_3 \hat{H}_1 + \mu \hat{L}_3$, where $\mu'^2 = \mu^2 + \epsilon_3^2$. The main feature of this basis is that BRpV is removed from the superpotential. Indeed, the superpotential in the rotated basis is given by

$$W = h_t \hat{Q}_3 \hat{U}_3 \hat{H}_2 + h_b \frac{\mu}{\mu'} \hat{Q}_3 \hat{D}_3 \hat{H}'_1 + h_\tau \hat{L}'_3 \hat{R}_3 \hat{H}'_1 - \mu' \hat{H}'_1 \hat{H}_2 + h_b \frac{\epsilon_3}{\mu'} \hat{Q}_3 \hat{D}_3 \hat{L}'_3, \quad (2)$$

with R–Parity non–conservation present in the form of TRpV, with an equivalent λ –coupling given by $\lambda'_{333} \equiv h_b \epsilon_3 / \mu'$.

Although BRpV is removed from the superpotential, it is reintroduced in the soft terms. The relevant terms in the original basis are

$$V_{soft} = m_{H_1}^2 |H_1|^2 + M_{L_3}^2 |\tilde{L}_3|^2 - [B \mu H_1 H_2 - B_3 \epsilon_3 \tilde{L}_3 H_2 + h.c.] + \dots \quad (3)$$

After performing the rotation described above, the soft lagrangian becomes

$$V_{soft} = m'^2_{H_1} |H'_1|^2 + M'^2_{L_3} |\tilde{L}'_3|^2 - [B' \mu' H'_1 H_2 - \frac{\epsilon_3 \mu}{\mu'^2} \Delta m^2 \tilde{L}'_3 H'_1 - \frac{\epsilon_3 \mu}{\mu'} \Delta B \tilde{L}'_3 H_2 + h.c.] + \dots \quad (4)$$

where we have defined $m'^2_{H_1} = (m_{H_1}^2 \mu^2 + M_{L_3}^2 \epsilon_3^2) / \mu'^2$ and $M'^2_{L_3} = (m_{H_1}^2 \epsilon_3^2 + M_{L_3}^2 \mu^2) / \mu'^2$ as the new scalar masses, and $B' = (B \mu^2 + B_3 \epsilon_3^2) / \mu'^2$ as the new B –term. The last two terms, where $\Delta m^2 \equiv m_{H_1}^2 - M_{L_3}^2$ and $\Delta B \equiv B_3 - B$, violate R–Parity bilinearly. It is clear that these two terms induce a non–zero vev for the rotated tau–sneutrino field $v'_3 \equiv \epsilon_3 v_1 + \mu v_3$. In models with universality of soft terms, the vev v'_3 is small because Δm^2 and ΔB are

radiatively generated at the weak scale and proportional to the bottom quark Yukawa coupling. In this case, using the tadpole equations [7], v'_3 can be approximated by

$$v'_3 \approx -\frac{\epsilon_3 \mu}{\mu'^2 m_{\tilde{\nu}_\tau}^2} (v'_1 \Delta m^2 + \mu' v_2 \Delta B) \quad (5)$$

where we have introduced

$$m_{\tilde{\nu}_\tau}^2 \equiv \frac{m_{H_1}^2 \epsilon_3^2 + M_{L_3}^2 \mu^2}{\mu'^2} + \frac{1}{8}(g^2 + g'^2)(v'_1)^2 - v_2^2. \quad (6)$$

This mass reduces to the tau–sneutrino mass of the MSSM in the $\epsilon_3 \rightarrow 0$ limit.

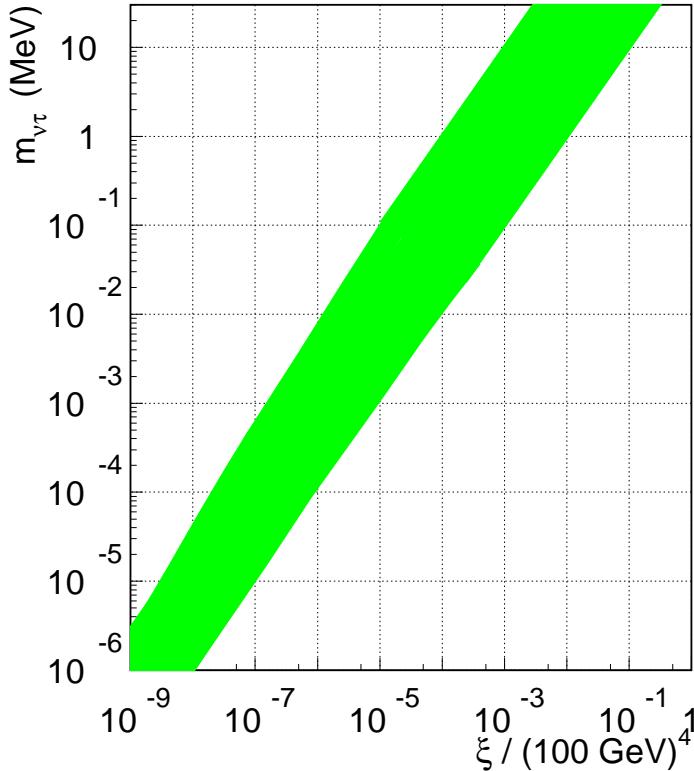


Figure 1: Tau–neutrino mass m_{ν_τ} as a function of the parameter $\xi \equiv v'_3 \mu'^2$, where v'_3 is the tau–sneutrino vacuum expectation value in the rotated basis.

The tau–neutrino acquires a mass because it mixes with the neutralinos. In a see-saw type of mechanism, with the neutralino masses playing the role of a high scale and v'_3 as the low scale, the tau–neutrino mass is approximately given by the expression

$$m_{\nu_\tau} \approx -\frac{(g^2 M' + g'^2 M) \mu'^2 v'_3}{4 M M' \mu'^2 - 2(g^2 M' + g'^2 M) v'_1 v_2 \mu'} \quad (7)$$

which is naturally small because of eq. (5). The tau–neutrino mass depends strongly on the tau–sneutrino vev v'_3 as it can be appreciated from Fig. 1. In this figure we plot m_{ν_τ}

as a function of the parameter $\xi \equiv (\epsilon_3 v_1 + \mu v_3)^2 = v_3'^2 \mu'^2$. We easily find solutions with neutrino masses from the collider limit of 17 MeV down to eV. The width of the band in Fig. 1 is related to the parameter $\tan \beta = v_2/v_1$, with the smaller (larger) values of $\tan \beta$ concentrated at the left (right) of the band.

2 Unconstrained MSSM–BRpV and $B(b \rightarrow s\gamma)$

By unconstrained MSSM–BRpV we understand the model where all soft parameters are independent at the weak scale, *i.e.*, not embedded into supergravity. We study the predictions of this model on the branching ratio $B(b \rightarrow s\gamma)$ varying randomly the soft parameters at the weak scale [8].

It is well known that the decay $b \rightarrow s\gamma$ is sensible to physics beyond the Standard Model (SM). The reason is that this decay is forbidden at tree level, and one-loop contributions from new physics compete with the SM contribution itself. The theoretical prediction of the decay rate in the SM, where W -bosons and top quarks contribute in the loops, is $B(b \rightarrow s\gamma) = (3.28 \pm 0.33) \times 10^{-4}$ [9] including NLO QCD corrections [9, 10, 11]. This prediction is in agreement at the 2σ level with the CLEO official measurement $B(b \rightarrow s\gamma) = (2.32 \pm 0.57 \pm 0.35) \times 10^{-4}$ [12]. Conservatively, this measurement implies $1.0 \times 10^{-4} < B(b \rightarrow s\gamma) < 4.2 \times 10^{-4}$ at 95% C.L., which has been modified by the preliminary measurement $2.0 \times 10^{-4} < B(b \rightarrow s\gamma) < 4.5 \times 10^{-4}$ at 95% C.L. reported in [13] after including more data.

In two Higgs doublets models of type II, where one Higgs doublet gives mass to the up-type fermions and the other to the down-type fermions, strong constraints on the charged Higgs mass are obtained because the H^\pm contribution adds to the W^\pm contribution to $B(b \rightarrow s\gamma)$ [14]. These strong constraints on m_{H^\pm} are also valid in supersymmetric (SUSY) models with large superpartners masses, although relaxed at high values of $\tan \beta$ [15] due to large corrections to the charged Higgs coupling to quarks [16].

In SUSY models with light superpartners the strong constraints on the charged Higgs mass are no longer valid after the inclusion of chargino, neutralino, and gluino loops along with squarks [17, 18, 19, 20, 21, 22]. This is specially due to the chargino loops which can cancel completely the charged Higgs loop. Supergravity models (SUGRA), with universality of soft mass parameters at the unification scale and with radiative breaking of the electroweak symmetry, are more constrained ruling out most of the parameter space if $\mu < 0$ and large $\tan \beta$ [22].

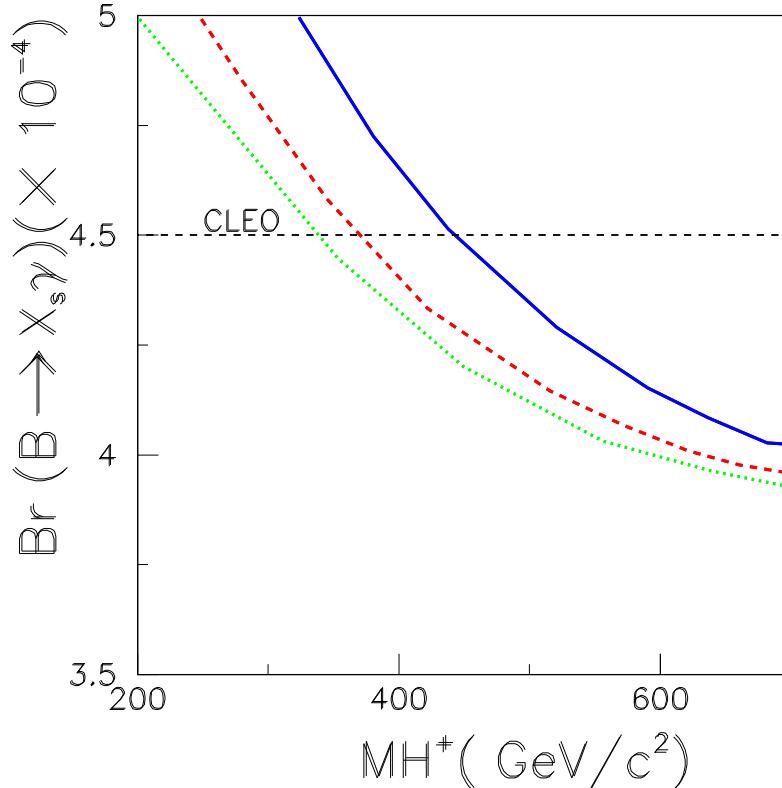


Figure 2: Lower limit of the branching ratio $B(b \rightarrow s\gamma)$ as a function of the charged Higgs mass m_{H^\pm} in the limit of very heavy squark masses. In solid is the MSSM and the other two curves correspond to BRpV. The horizontal line corresponds to the upper experimental limit from CLEO.

It has been shown that Bilinear R-Parity Violation can relax the bounds on the charged Higgs mass [8]. In this model new particles contribute in the loops to $B(b \rightarrow s\gamma)$. Charginos mix with the tau lepton ², therefore, the tau lepton contribute to the decay rate together with up-type squarks in the loops. Nevertheless, it was demonstrated that the tau contribution can be neglected [8]. In a similar way, the charged Higgs boson mixes with the two staus [24] forming a set of four charged scalars, one of them being the charged Goldstone boson. In this way, the staus contribute to the decay rate together with up-type quarks in the loops.

The four charged scalars in the original basis are $\Phi^\pm = (H_1^\pm, H_2^\pm, \tilde{\tau}_L^\pm, \tilde{\tau}_R^\pm)$ and the corresponding mass matrix is diagonalized after the rotation $\mathbf{S}^\pm = \mathbf{R}_{S^\pm} \Phi^\pm$ where \mathbf{S}_i^\pm , $i = 1, 2, 3, 4$ are the mass eigenstates (one of them the unphysical Goldstone boson). One of the massive charged scalars has similar properties to the charged Higgs of the MSSM.

² This mixing is not in conflict with the well measured tau couplings to gauge bosons [23].

In BRpV we call the “charged Higgs boson” to the charged scalar whose couplings to quarks are larger, *i.e.*, maximum $(\mathbf{R}_{S^\pm}^{i1})^2 + (\mathbf{R}_{S^\pm}^{i2})^2$. Nevertheless, for comparison we have also study the case in which the “charged Higgs boson” corresponds to the charged scalar with largest components to the rotated Higgs fields $H_1'^\pm$ and $H_2'^\pm$, *i.e.*, maximum $(\mathbf{R}'_{S^\pm}^{i1})^2 + (\mathbf{R}'_{S^\pm}^{i2})^2$.

We neglect in this calculation the contribution of neutralinos, because it is small [17], and that of the gluino whose different squark contributions tend to cancel with each other [22]. In addition, if gaugino masses are universal at the GUT scale, gluinos must be rather heavy considering the bound on the chargino mass from LEP2 [25], which makes the contribution smaller. We ignore the light gluino window [26] because it is inconsistent with the experimental bound on the mass of the lightest Higgs boson in the MSSM [27].

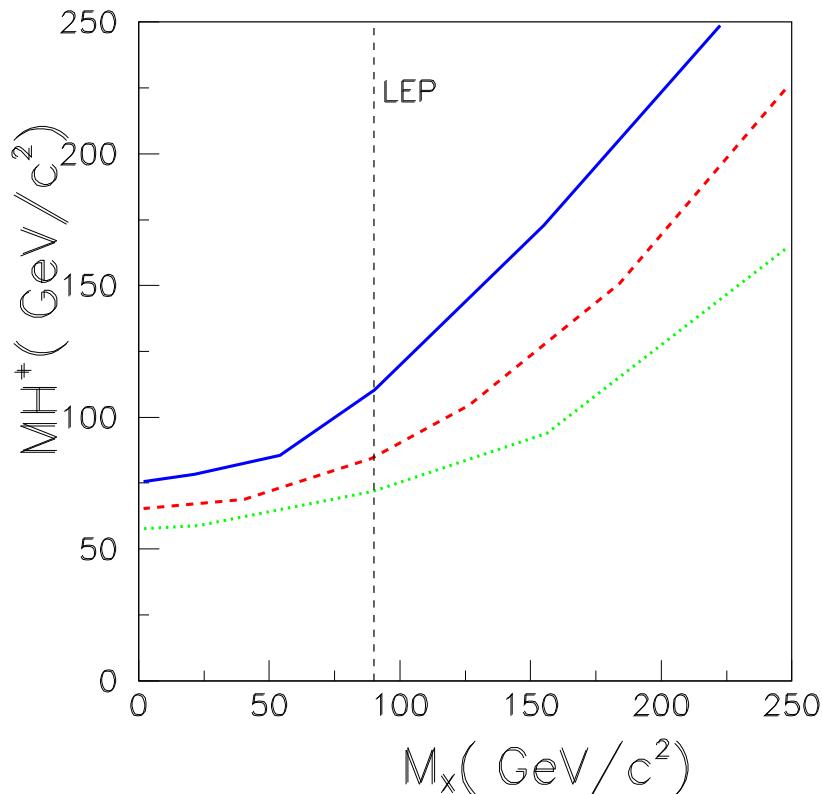


Figure 3: Lower limit of the charged Higgs mass m_{H^\pm} as a function of the lightest chargino mass $m_{\chi_1^\pm}$ compatible with the CLEO measurement. In solid is the MSSM and the other two curves correspond to BRpV. The vertical line is the experimental lower limit on $m_{\chi_1^\pm}$ from LEP.

In the limit of very heavy squarks, the strong constraints imposed on the charged Higgs mass of the MSSM are relaxed in the MSSM–BRpV as can be appreciated in Fig. 2.

Above and to the right of the solid line are the solutions of the MSSM consistent with the CLEO measurement of $B(b \rightarrow s\gamma)$. Without considering theoretical uncertainties, the limit on the charged Higgs mass is $m_{H^\pm} > 440$ GeV. This bound is relaxed by about 70 to 100 GeV in BRpV as can be seen from the dotted and dashed lines. If a 10% theoretical uncertainty is considered, the MSSM bound reduces to $m_{H^\pm} > 320$ GeV, but the BRpV bound decreased as well such that the reduction of the bound is maintained. The dotted line corresponds to the charged Higgs with largest couplings to quarks, a definition that makes more sense in our calculation. The dashed line corresponds to a charged Higgs defined by maximum component along $H_1'^\pm$ and $H_2'^\pm$.

Another interesting case is the region of parameter space where the charged Higgs and the charginos are light. The limits on light m_{H^\pm} in the MSSM are also relaxed in BRpV as shown in Fig. 3. Solutions consistent with the CLEO measurement of $B(b \rightarrow s\gamma)$ in the MSSM lie over and to the left of the solid curve, implying that a chargino heavier than 90 GeV requires a charged Higgs heavier than 110 GeV. This bound on m_{H^\pm} is relaxed in BRpV by 25 to 35 GeV as showed by the other two curves. In particular, the charged Higgs can be lighter than m_W and observable at LEPII.

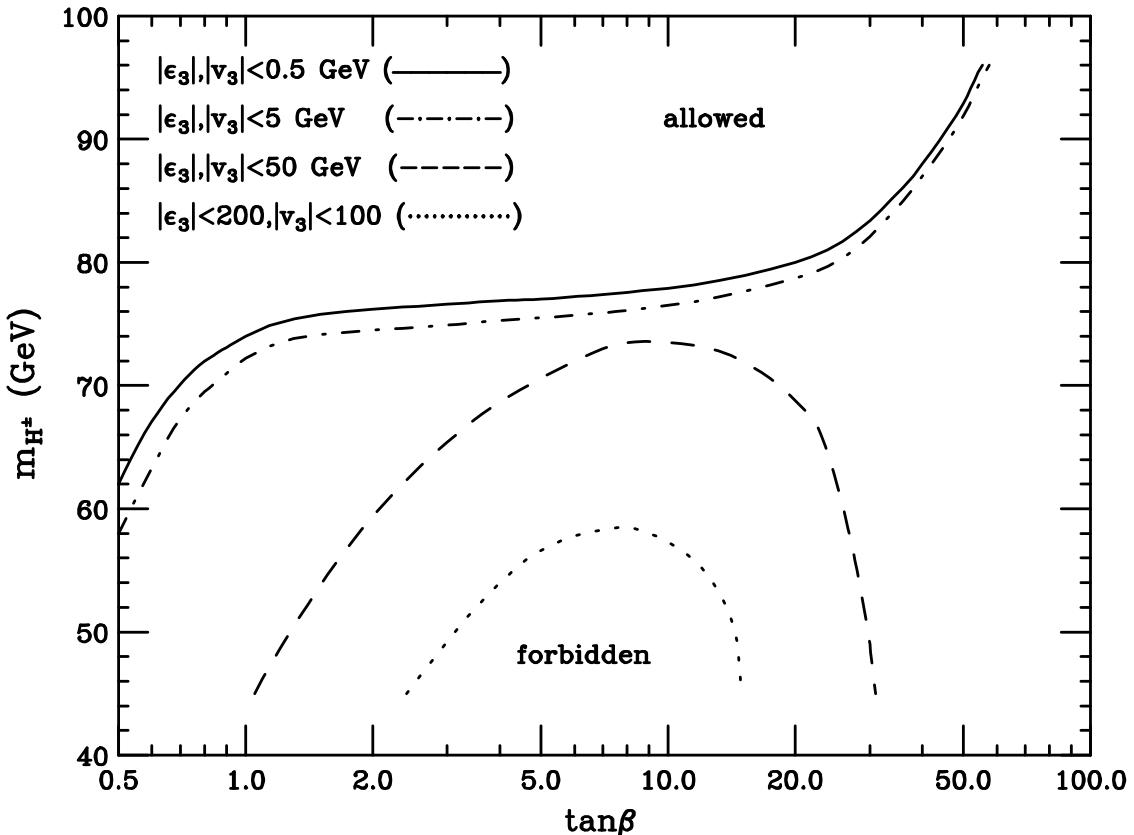


Figure 4: Lower limit of the charged Higgs mass m_{H^\pm} as a function of the

In the MSSM, the charged Higgs can be lighter than m_W after the inclusion of

radiative corrections [28] in some corners of parameter space. In BRpV this situation is not so rare [24]. In Fig. 4 we show that for moderate values of ϵ_3 the charged Higgs mass may be lower than m_W . Nevertheless, to the normal MSSM decay modes we need to add the R-Parity violating decay modes $H^\pm \rightarrow \tau^\pm \tilde{\chi}_1^0$ and $H^\pm \rightarrow \tilde{\chi}_1^\pm \nu_\tau$ which can be dominant at low $\tan \beta$ [24].

3 BRpV Embedded into Supergravity and the Lightest Higgs Mass

The BRpV model can be successfully [7] embedded into SUGRA with radiative breaking of the electroweak symmetry [29] and universality of soft masses. The electroweak symmetry is broken through the vacuum expectation value (vev) of the tau-sneutrino, in addition to the two Higgs field vevs, and it contributes to the mass of the gauge bosons. The correct vev's are found by imposing the three tadpole equations, where one-loop tadpoles corrections are important for the tau-sneutrino as well as for the two Higgs fields [7].

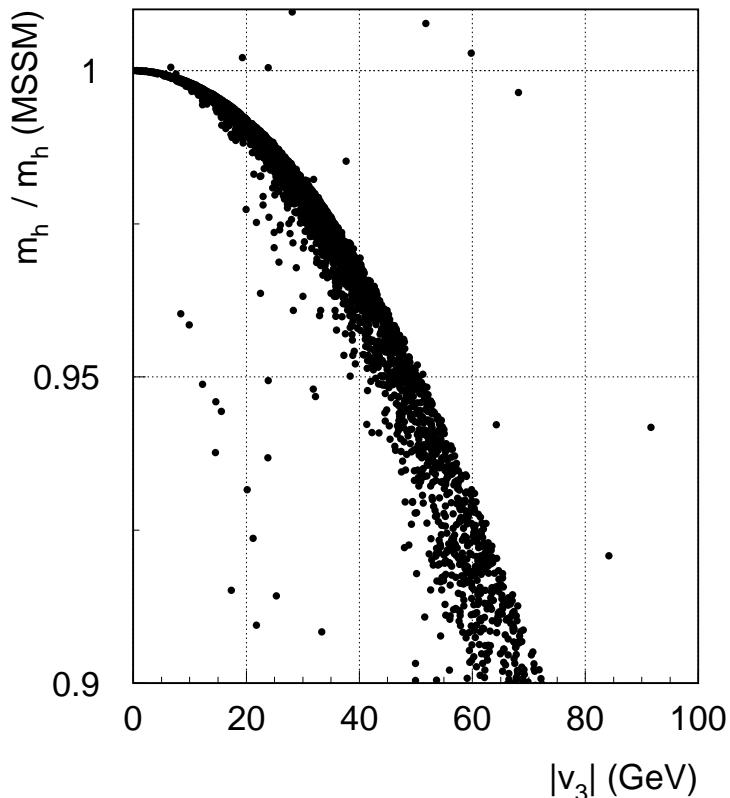


Figure 5: Ratio between the lightest Higgs boson mass m_h in BRpV and the same mass in the MSSM ($v_3 = 0$ limit) as a function of the tau-sneutrino vacuum expectation value.

In MSSM–BRpV the CP–even Higgs bosons mix with the real part of the tau–sneutrino [3]. The effect of this mixing is to lower the mass of the lightest scalar as can be appreciated in Fig. 5, with the exception of a few, statistically insignificant, exceptional points.

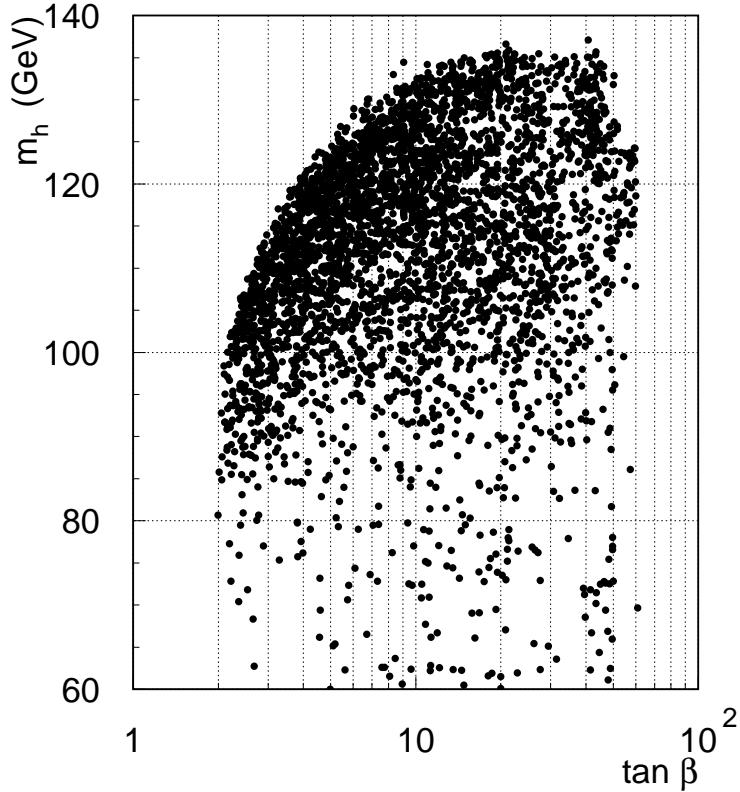


Figure 6: Lightest neutral Higgs mass in BRpV as a function of $\tan \beta$. The lower bound on $\tan \beta$ is due to the non-perturbativity of the top quark Yukawa coupling.

In Fig. 6 we see that the upper bound on the lightest CP–even Higgs mass does not change neither. In this figure we have included in the 3×3 Higgs mass matrix only the radiative corrections proportional to the fourth power of the top quark mass [30]. The neglected corrections brings down the upper bound in several GeV. As in the MSSM, the upper bound for m_h decreases when $\tan \beta$ decreases, reaching in BRpV somewhat less than 100 GeV if $\tan \beta = 2$. This implies that experimental searches for h at CERN also tests this model if $\tan \beta$ is close to unity, as it occurs in the MSSM ³.

It is interesting to note that, since the vacuum stability bound on the SM Higgs boson mass is about 135 GeV [31], the measurement of the Higgs boson mass can distinguish

³ Simple formulas for the one-loop radiatively corrected lightest neutral Higgs mass m_h when $\tan \beta$ is close to one can be found in [30].

between the BRpV model and the SM (with no physics below $\sim 10^{10}$ GeV) in the same way as it can distinguish between the MSSM (or the NMSSM) and the SM [32].

4 Unification of Yukawa Couplings

In the MSSM, bottom–tau Yukawa unification is achieved at two disconnected regions at low and high values of $\tan\beta$ [33]. This can be appreciated in the “inverted U” shaped region in the plane m_t – $\tan\beta$ shown in Fig. 7. The two horizontal lines correspond to the 1σ determination of the top quark mass [34] at Fermilab ⁴.

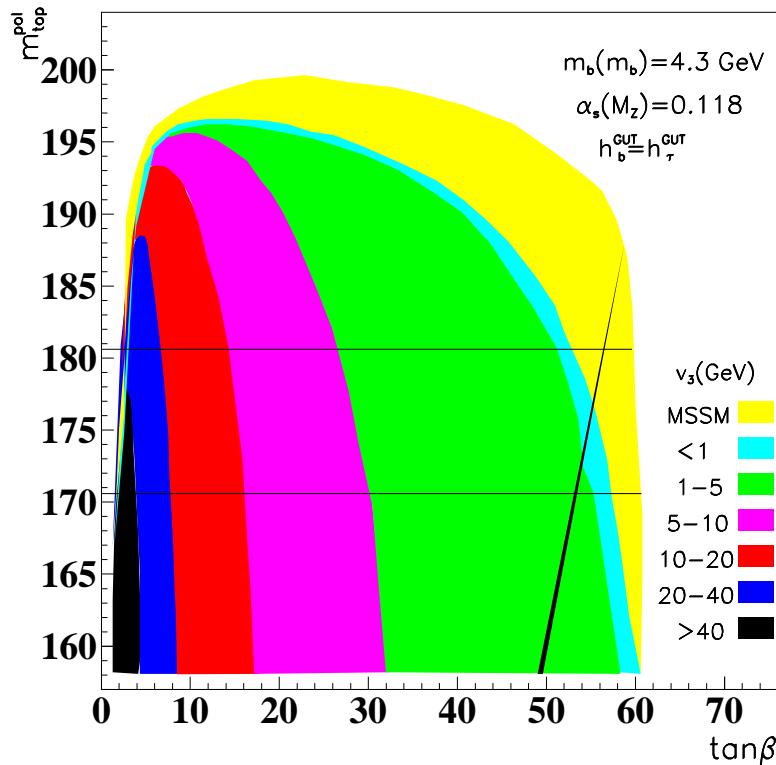


Figure 7: Regions of the m_t - $\tan\beta$ plane where bottom–tau Yukawa coupling can be achieved in the MSSM and in BRpV. In the case of BRpV regions are labeled by the tau–sneutrino vev v_3 . The inclined straight line corresponds to top–bottom–tau Yukawa unification.

In BRpV the unification of couplings is modified [36] in an important way (see [37] for the TRpV case). In Fig. 7 we see that by choosing the value of the tau–sneutrino

⁴ We note that the new decay modes of the top quark present in the BRpV model impose only mild constraints on the BRpV parameter ϵ_3 [35].

vev v_3 we can achieve bottom–tau Yukawa unification at any value of $\tan \beta$ provided we keep inside the perturbativity region $2 \lesssim \tan \beta \lesssim 60$. The high $\tan \beta$ region where top–bottom–tau unification is found is twice as big in BRpV compared with the MSSM. The $t - b - \tau$ unification is found at values of $v_3 < 5$ GeV, therefore, it would rule out regions of parameter space where the bilinear violation of R–Parity is large. We note that in Fig. 7 we define $\tan \beta = v_2/v_1$ to preserve the MSSM definition. Another possibility is to define $\tan \beta' = v_2/\sqrt{v_1^2 + v_3^2}$ which has the advantage of being invariant under rotations on the $L_3 - H_1$ plane. We have checked that Fig. 7 does not change appreciably when plotted against $\tan \beta'$.

The reason why $b - \tau$ unification in BRpV fills the intermediate regions of $\tan \beta$ can be understood as follows. First of all, we notice that the quark and lepton masses are related to the different vevs and Yukawa couplings in the following way

$$m_t^2 = \frac{1}{2} h_t^2 v_2^2, \quad m_b^2 = \frac{1}{2} h_b^2 v_1^2, \quad m_\tau^2 = \frac{1}{2} h_\tau^2 v_1^2 (1 + \delta), \quad (8)$$

where δ depends on the parameters of the chargino/tau mass matrix and is positive [24, 36]. This implies that the ratio of the bottom and tau Yukawa couplings at the weak scale is given by

$$\frac{h_b}{h_\tau}(m_{weak}) = \frac{m_b}{m_\tau} \sqrt{1 + \delta} \quad (9)$$

and grows as $|v_3|$ is increased.

On the other hand, if h_b and h_τ unify at the GUT scale, then at the weak scale its ratio can be approximated by

$$\frac{h_b}{h_\tau}(m_{weak}) \approx \exp \left[\frac{1}{16\pi^2} \left(\frac{16}{3} g_s^2 - 3h_b^2 - h_t^2 \right) \ln \frac{M_{GUT}}{m_{weak}} \right] \quad (10)$$

implying that the combination $3h_b^2 + h_t^2$ should decrease when $|v_3|$ increases.

In the MSSM region of high $\tan \beta$ the bottom quark Yukawa coupling dominates over the top one, and the opposite happens in the region of low $\tan \beta$. Therefore, at high (low) values of $\tan \beta$, the Yukawa coupling h_b (h_t) will decrease if $|v_3|$ increases, which implies an increase of v_1 (v_2) in order to keep constant the quark masses. Similarly, in order to keep constant the W mass, $m_W^2 = \frac{1}{4} g^2 (v_1^2 + v_2^2 + v_3^2)$, the vev v_2 (v_1) decreases at the same time. This implies that unification occur at lower (higher) values of $\tan \beta$ as $|v_3|$ increases, explaining what we see in Fig 7.

The fact that bottom–tau unification occurs at any value of $\tan \beta$ in BRpV seems important to us considering that in the MSSM the low $\tan \beta$ region is disfavoured by the non observation of the lightest Higgs boson. In addition, the high $\tan \beta$ region is dis-

favoured because it is usually difficult to find the correct electroweak symmetry breaking. These difficulties are avoided in BRpV if the sneutrino vev is sufficiently large.

5 Conclusions

In its simplest form, Bilinear R–Parity Violation is a one parameter extension of the MSSM which can be successfully embedded into SUGRA with radiative breaking of the electroweak symmetry and universality of soft masses. This is achieved through the running of the same RGEs of the MSSM since BRpV does not introduce new interactions. Therefore it is a very simple framework to study R–Parity violating phenomena. In addition, BRpV generates a tau–neutrino mass which, in models with universality of soft masses, is radiatively generated and proportional to the bottom quark Yukawa coupling squared, therefore, naturally small.

In BRpV charged Higgs bosons mix with the staus, and because of this, staus contribute to the decay $b \rightarrow s\gamma$. In an unconstrained version of the model we have showed that the bounds on the charged Higgs boson mass from $B(b \rightarrow s\gamma)$ are relaxed by ~ 100 GeV in the heavy squark limit, and by ~ 30 GeV in the light chargino and light charged Higgs limit. In this case, charged Higgs lighter than the W –gauge boson are possible and observable at LEP2. Nevertheless, R–Parity violating decay modes will compete with the traditional decay modes of the charged Higgs in the MSSM. In a similar way, the neutral CP–even Higgs bosons mix with the real part of the tau–sneutrino. In general, this mixing lowers the Higgs mass but leaves the upper bound unchanged.

Finally, we have shown that it is much easier to find unification of the bottom and tau Yukawa couplings in BRpV than in the MSSM. By choosing the value of the tau–sneutrino vacuum expectation value, $b - \tau$ unification can be achieved at any value of $\tan \beta$. Unification of $t - b - \tau$ can be found at high values of $\tan \beta$ but in a region twice as large than in the MSSM.

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